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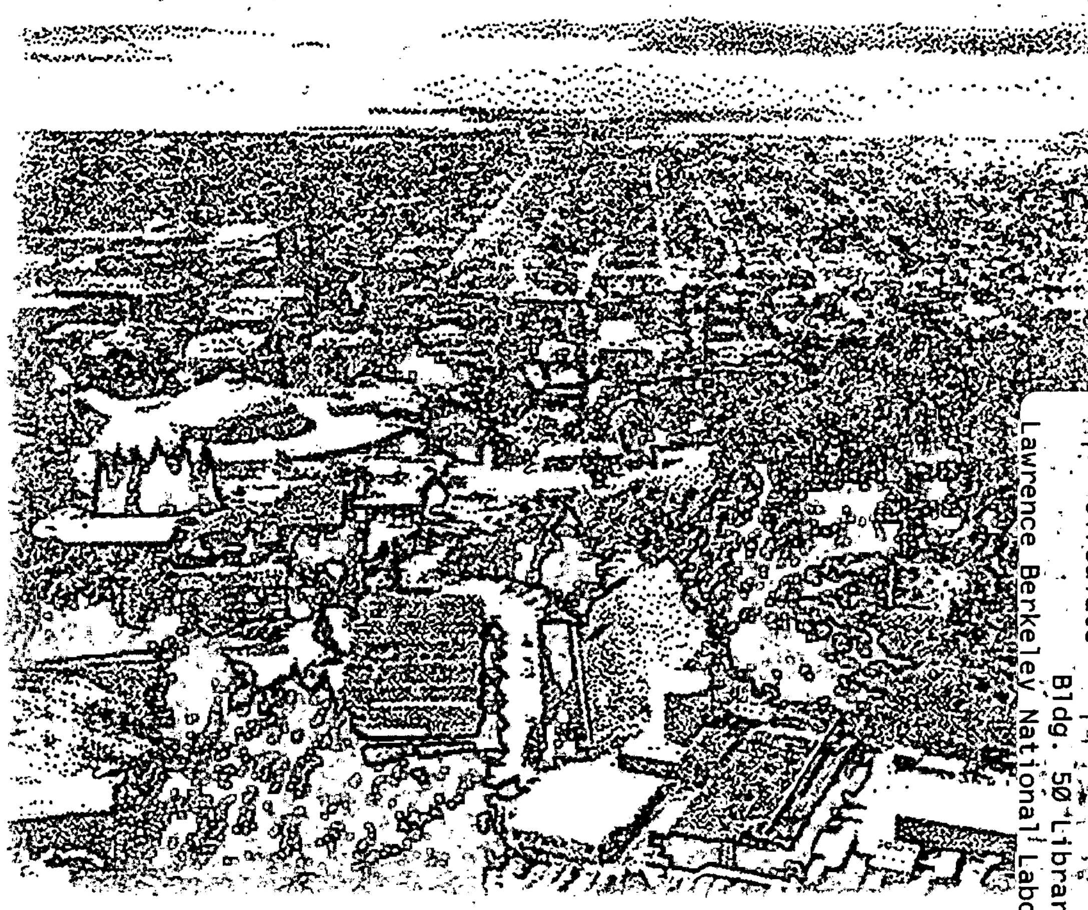

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Regulating the Baryon Asymmetry in No-Scale Affleck-Dine Baryogenesis *

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Abstract

In supergravity models (such as standard superstring constructions) that possess a Heisenberg symmetry, supersymmetry breaking by the inflationary vacuum energy does not lift flat directions at tree level. One-loop corrections give small squared masses that are negative ($\sim -g^2H^2/(4\pi)^2$) for all flat directions that do not involve the stop. After inflation, these flat directions generate a large baryon asymmetry; typically $n_B/s \sim$ O(1). We consider mechanisms for suppressing this asymmetry to the observed level. These include dilution from inflaton or moduli decay, GUT nonflatness of the *vev* direction, and higher dimensional operators in both GUT models and the MSSM. We find that the observed BAU can easily be generated when one or more of these effects is present.

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1 Introduction

Supersymmetric (SUSY) theories possess a unique and efficient mechanism for the generation of a cosmological baryon asymmetry through the decay of scalar fields along nearly F- and D-flat directions of the scalar potential, using baryon number violation arising in GUTS [1], or other sources of baryon or lepton number violation such as neutrino mass effects [2, 3, 4]. In its original form, applied to SUSY-GUTS, this Affleck-Dine (AD) mechanism not only produces a baryon asymmetry but also entropy from the decaying fields and results in a net baryon-to-entropy ratio, $n_B/s \sim$ O(1) [1, 5]. In fact, one of the main problems associated with this mechanism in the GUT case is the dilution of the asymmetry down to acceptable levels of order 10^{-11}–10^{-10} [6]. In the context of GUTS after inflation, the AD scenario can be shown [7] to produce an asymmetry of the desired magnitude. During inflation, scalar fields with masses, $m < H$, are driven by quantum fluctuations to large vacuum values (but less than M_P) which become the source of the scalar field oscillations once inflation is over: however, inflation supplies an additional source of entropy from inflaton decays.

Recently, it has been argued that the simple picture of driving scalar fields to large vacuum values along flat directions during inflation is dramatically altered in the context of supergravity [8]. During inflation, the Universe is dominated by the vacuum energy density, $V \sim H^2 M_P^2$. The presence of a nonvanishing and positive vacuum energy density indicates that supersymmetry is broken and soft masses of order of $\sim H$ are generated [9]. If such masses are generated along otherwise flat directions, the AD mechanism would be inoperative. While this is true for minimal supergravity, it was shown [10] that in no-scale supergravity [11], or more generally in any supergravity theory with a Heisenberg symmetry of the kinetic function [12], such corrections are absent at the tree level. It was also shown [10] that, at the one-loop level, corrections to the scalar potential are generated along flat directions and, for all directions which do not involve stops, these corrections to the mass squared are negative of magnitude $\sim 10^{-2} H^2$. This is large enough to insure that the *vev*'s along flat directions do indeed run to $\phi_0 \sim M_P$, resulting once more in a baryon asymmetry of order unity.

In this letter, we will systematically consider possible effects which may contribute to diluting a baryon asymmetry $n_B/s \sim$ O(1) from the AD mechanism. We will consider in turn the effects of dilution by an inflaton and/or a massive moduli field, the effects of higher dimensional operators, and the effects of an intermediate scale. We find that the dilution by an inflaton or moduli is insufficient in generic supergravity models, but may be sufficient in the context of certain string-derived models. In the absence of these effects the presence of the first nonrenormalizable contribution to the superpotential (of dimension 4), suppressed by the GUT scale, can lead to the desired asymmetry.

We also consider baryogenesis in the context of the minimal supersymmetric standard model, with only the baryon number violation, and nonperturbative corrections to the superpotential, implied by quantum gravity effects at the Planck scale. We find that these effects are sufficient to generate the observed baryon asymmetry of the universe (BAU) in the case of the minimal supersymmetric standard model with standard 4-dimensional (super)gravity. We also consider the possibility that the asymmetry is diluted by sphaleron effects at the electroweak scale.

2 Generalities

The contribution to scalar masses can be derived from the scalar potential in a supergravity model described by a Kähler potential G [13],

$$V = e^G \left[G_i (G^{-1})^i_j G^j - 3 \right] \tag{1}$$

where $G_i = \partial G / \partial \phi^i$ and $G^i = \partial G / \partial \phi^*_i$ and we are setting the D-terms equal to 0, which is true for gauge singlets and along the D-flat directions. In minimal supergravity, the Kähler potential is defined by

$$G = zz^* + \phi^*_i \phi^i + \ln |\overline{W}(z) + W(\phi)|^2 \tag{2}$$

where z is a Polonyi-like field [14] needed to break supersymmetry, and we denote the scalar components of the usual matter chiral supermultiplets by ϕ^i. W and $\overline{W}$ are the superpotentials of ϕ^i and z respectively. In this case, the scalar potential becomes

$$V = e^{zz^* + \phi^*_i \phi^i} \left[|\overline{W}_z + z^*(\overline{W} + W)|^2 + |W_{\phi^i} + \phi^*_i (\overline{W} + W)|^2 - 3|(\overline{W} + W)|^2 \right]. \tag{3}$$

In the above expression for V, one finds a mass term for the matter fields ϕ^i, $e^G \phi^*_i \phi^i = m^2_{3/2} \phi^*_i \phi^i$ [15]. Scalar squared masses will pick up a contribution of order $m^2_{3/2} \sim V/M_P^2 \sim H^2$, thus lifting the flat directions and potentially preventing the realization of the AD scenario as argued in [8]. Below we use reduced Planck units: $M_P = 1/\sqrt{8\pi G_N} = 1$. All flat directions in the matter sector are lifted as well.

In models with a "Heisenberg symmetry" [12] the dangerous cross-term is absent. If in addition to the field z and matter fields ϕ^i, we allow for hidden sector fields y^i, the

"Heisenberg symmetry" is defined by

$$\delta z = \epsilon_i^* \phi^i, \qquad \delta\phi^i = \epsilon^i, \qquad \delta y^a = 0, \tag{4}$$

where ϵ^i are complex parameters and ϵ_i^* their complex conjugates. The combination

$$\eta \equiv z + z^* - \phi_i^* \phi^i. \tag{5}$$

and y_i are invariant. Let us assume this is a symmetry of the kinetic function in the Kähler potential. We also require that the field z does not have a coupling in the superpotential. Then the most general Kähler potential becomes

$$G = f(\eta) + \ln |W(\phi)|^2 + g(y), \tag{6}$$

where the superpotential W is a holomorphic function of ϕ^i only, and the fields η and ϕ^i are regarded as independent degrees of freedom [16, 17, 10], yielding a potential [18]:

$$V = e^{f(\eta)+g(y)} \left[\left(\frac{f'^2}{f''} - 3 \right) |W|^2 - \frac{1}{f'} |W_i|^2 + g_a (g^{-1})^a_b g^b |W|^2 \right]. \tag{7}$$

It is important to notice that the cross term $|\phi_i^* W|^2$ has disappeared in the scalar potential. Because of the absence of the cross term, flat directions remain flat even during inflation. A detailed discussion of this result is found in [10].

The above result however, is only valid at the tree-level. Since gravitational interactions preserve the Heisenberg symmetry at one-loop [12], the only possible contribution to the mass of the flat directions can come from either gauge interactions or superpotential couplings that contribute to the renormalized Kähler potential. This too was computed in [10] from the general results [19, 20] for the one-loop corrected supergravity Lagrangian. Assuming that inflation is driven by an F-term rather than a D-term (the result is similar if $\langle D \rangle \neq 0$), the one-loop corrected mass is

$$\begin{aligned} \left(m^2\right)^j_i &= \frac{\ln(\Lambda^2/\mu^2)}{32\pi^2} \Bigg\{ h_{ikl} h^{*jkl} \left[\alpha\langle V \rangle + m_{3/2}^2 \left(5\frac{f'^2}{f''} + 2\frac{f' f'''}{f''^2} - 10 - \frac{f'^4}{f''^2} \right) \right] \\ &\quad - 4\delta^j_i g_a^2 C_2^a(R_i) \left[\beta\langle V \rangle + m_{3/2}^2 \left(\frac{f' f'''}{f''^2} - 2 \right) \right] \Bigg\}, \end{aligned} \tag{8}$$

where the vacuum energy $\langle V \rangle$ and the gravitino mass $m_{3/2}$ are their values during inflation. We are assuming that the inflaton, ψ, is one of the ϕ^i or a hidden sector y^i field. $\mu \geq \langle V \rangle$

is the appropriate infrared cut-off in the loop integral, and Λ is the cutoff scale below which the effective supergravity Lagrangian given by the Kähler potential eq. (6) is valid; $\Lambda = 1$ in many models but can be lower if the effective potential for inflation is generated at an intermediate scale. The h's are Yukawa coupling constants $g_a, C_2^a(R_i)$ are the coupling constant and matter Casimir for the factor gauge group G_a, and the parameter α, β, are model dependent. In the no-scale case $f = -3\ln\eta$, the result (8) reduces to

$$\left(m^2\right)_i^j = \frac{\ln(\Lambda^2/\mu^2)}{32\pi^2}\langle V\rangle\left[\alpha h_{ikl}h^{*jkl} - 4\beta\delta_i^j g_a^2 C_2^a(R_i)\right]. \qquad (9)$$

If the inflaton is one of the ϕ^i ($\phi_0 \neq 0$), $\alpha = \beta = \frac{2}{3}$. If the inflaton is in the hidden sector ($y_a \neq 0$), $\alpha < 0$, $\beta = 1$ if the inflaton is not the dilaton s : $\langle s\rangle = g^{-2}$. If the inflaton is the standard string dilaton, $\beta = \alpha = -1$. In all but the last case the masses are negative if gauge couplings dominate Yukawa couplings. We will assume that the inflaton is not the dilaton.

The vacuum energy in Eq. (9) is related to the expansion rate by $H^2 = (1/3)\langle V\rangle$, and hence the typical mass of the flat directions is $m^2 \simeq 10^{-2}H^2$ during inflation. However, from Eq. (9), we see that for all scalar matter fields aside from stops, the contribution to the mass squared is *negative* as the Yukawa couplings are smaller than the gauge couplings. Thus any flat direction not involving stops will have a negative contribution at one-loop without an ad hoc choice of the parameters. Now, as argued in [10], even though fluctuations will begin the growth of ϕ_0, the classical equations of motion soon take over. The classical equations of motion drive ϕ_0 as $(-m^2)t$ which is smaller than the quantum growth only for $Ht < H^4/m^4$. Then for $Ht > H^2/(-m^2)$, the classical growth of ϕ_0, becomes nonlinear $\sim He^{-m^2t/3H}$, and ϕ_0 will run off to its minimum determined by the one-loop corrections to ϕ^4, which are again of order V. An explicit one-loop calculation [19, 20] shows that the effective potential along the flat direction has the form

$$V_{eff} \simeq \frac{g^2}{(4\pi)^2}\langle V\rangle\left(-2\phi^2\log\left(\frac{\Lambda^2}{g^2\phi^2}\right) + \phi^2\right) + \mathcal{O}(\langle V\rangle)^2, \qquad (10)$$

where Λ is the cutoff of the effective supergravity theory, and has a minimum around $\phi \simeq 0.5\Lambda$. This is consistent because this effective potential is only of order $-\langle V\rangle g^2/(4\pi)^2$ and is a small correction to the inflaton energy density which drives inflation. Thus, $\phi_0 \sim M_P$ will be generated and in this case the subsequent sfermion oscillations will dominate the energy density and a baryon asymmetry will result which is independent of inflationary parameters as originally discussed in [1, 5] and will produce $n_B/s \sim O(1)$.

For a ϕ_0 *vev* as large as M_P, coherent oscillations of the flat directions will persist long after the inflatons have decayed and the radiation produced in their decays has redshifted away. To see this [7], consider the energy density stored in the inflaton and the scalar fields making up our flat direction

$$\rho_\psi = m_\psi^2 \psi^2 \tag{11}$$

$$\rho_\phi = m_\phi^2 \phi^2 \tag{12}$$

We will assume that the initial *vev* for ψ is M_P and the initial *vev* for ϕ is ϕ_0, which as we saw above, due to the one-loop correction to the flat direction is also M_P. After inflation (the de Sitter expansion), the corrections in Eq. (10) are turned off and the flat direction is restored up to supersymmetry breaking effects of order $\tilde{m} \sim 10^{-16}$. If we consider only a generic model of inflation with one single input scale determined by the COBE normalization of the microwave background anisotropy, then $m_\psi \sim 10^{-7}$ [7, 3] and the inflatons begin oscillating about the minimum of the inflationary potential when the Hubble parameter $H \simeq m_\psi \psi / M_P \sim m_\psi$. Let us denote the value of the scale factor R at this point R_ψ. For $R > R_\psi$, the Universe expands as if it were matter dominated so that $H \simeq m_\psi (R_\psi/R)^{3/2}$. The two field system involving the flat direction and the inflaton was worked out in detail in [7]. We only quote the relevant results here. Inflatons decay at $R_{d\psi} \simeq (M_P/m_\psi)^{4/3} R_\psi$ when $\Gamma_\psi = m_\psi^3/M_P^2 \simeq H$ (their decay rate is assumed to be gravitational). After inflaton decay, the Universe is radiation dominated and $\rho_{r\psi} \simeq m_\psi^{2/3} M_P^{10/3} (R_\psi/R)^4$, where $\rho_{r\psi}$ is the energy density in the radiation produced by inflaton decay.

Similarly, the fields ϕ begin oscillating about the origin, when $H \sim \tilde{m} > \Gamma_\psi$. After ϕ's begin oscillating at $R = R_\phi$, their energy density is just $\rho_\phi \simeq \tilde{m}^2 \phi_0^2 (R_\phi/R)^3 \simeq m_\psi^2 \phi_0^2 (R_\psi/R)^3$. In the AD mechanism, one utilizes the flat directions lifted only by supersymmetry breaking effects

$$\int d^4\theta \zeta^* \zeta \frac{\phi^* \phi^* \phi \phi}{M_X^2} = \tilde{m}^2 \frac{\phi^* \phi^* \phi \phi}{M_X^2}, \tag{13}$$

where $\zeta = \tilde{m}\theta^2$ is the supersymmetry-breaking spurion. Such operators can carry non-vanishing baryon-numbers. The flat directions associated with baryon number violating operators store a net baryon number density $n_B = \epsilon(\phi_0^2/M_X^2)\rho_\phi/\tilde{m}$, where M_X is the scale of the baryon number violating operator, $M_X \sim 10^{-3}$ for GUT scale baryon number violation and ϵ is measure of CP violation in the oscillations and should be O(1). Oscillations along the flat direction decay when their decay rate $\Gamma_\phi \simeq \tilde{m}^3/\phi^2 \sim H$. If

$$\phi_0 > \tilde{m}^{5/12} M_P^{4/3}/m_\psi^{3/4} \sim 10^{-3/2} \tag{14}$$

the flat direction fields decay so late that the energy density in radiation from inflaton decays, $\rho_{r\psi}$ has redshifted away and the Universe is dominated by ϕ's when they decay. In this case, the resulting entropy in the Universe is produced by ϕ decay and $n_B/s \sim$ O(1) [1, 5]. (In fact, if the asymmetry were computed solely from sfermion decays as in [1], it might appear that a baryon to entropy ratio greater than 1, $n_B/s \sim (M_P/\tilde{m})^{1/6}$, were possible. However as argued by Linde [5], the maximum asymmetry achieved is ~ 1, independent of the the initial value of ϕ_0.) For $\phi_0 < \tilde{m}^{5/12} M_P^{4/3}/m_\psi^{3/4}$, ϕ's decay while the Universe is dominated by $\rho_{r\psi}$. The Universe will remain radiation dominated and the baryon-to entropy ratio is that computed in [7]

$$\frac{n_B}{s} \sim \frac{\epsilon {\phi_0}^4 {m_\psi}^{3/2}}{({M_X}^2 + \phi_0^2){M_P}^{5/2}\tilde{m}} \lesssim \min(1, 10^6(\frac{\phi_0}{M_P})^2, 10^{12}(\frac{\phi_0}{M_P})^4). \tag{15}$$

In eq. (15), when the condition (14) is satisfied, the asymmetry is O(1). For smaller ϕ_0, the asymmetry is given by one of the two expressions on the right of (15), depending on whether ϕ_0 is greater or less than M_X. Finally, for $\phi_0 < \tilde{m}^{5/2} M_P^3/m_\psi^{9/2}$, the flat direction fields decay before the inflatons, but in this case, too small an asymmetry would result. These results are distinguished along the right axes in the Figure. Values of $\tilde{m}$ and M_X were taken to be 10^{-16} and 10^{-3} respectively in eq. (15) and in the Figure. In the absence of the types of effects we discuss below, we would expect $\phi_0 \sim$ O(1) due to the one-loop potential (10) and our baryon asymmetry is too large. This is the source of our problem.

3 Dilution by Moduli Decay

Polonyi-like moduli (which we will label as y) have been known to be a source of embarrassment primarily due to their excessive entropy production [21]. In a conventional GUT baryogenesis scenario, the late decay of moduli would lead to an entropy dilution factor of O(10^{16}), and lead to a low reheat temperature of about 1 keV unless $m_y \gtrsim 10$ TeV, similar to the cosmological gravitino problem [22]. Furthermore, in models with R-parity conservation, it was pointed out that unless $m_y \gtrsim 10^6$ GeV, the decay of y oscillations will lead to an unacceptably large LSP relic density [23]. This is the shaded region on the left in the figure. Because of the potential problem with entropy production, the y's seem to be a natural place to look for the dilution of the AD baryon asymmetry.

For now, we will treat the moduli mass as a free parameter, m_y. Unless, $m_y < m_\psi$, the moduli would dominate the energy density of the Universe early on ($\rho_y \simeq m_y^2 y_0^2$ and $y_0 \sim M_P$) contrary to our assumption that inflation is driven by the inflaton. For $m_y < m_\psi$, the moduli oscillations live on after inflaton decay, and can subsequently come to dominate the energy density of the Universe. As in the ϕ-ψ system, when $\phi_0 > \tilde{m}^{5/12} M_P^{4/3} / m_y^{3/4}$, ϕ's come to dominate and a baryon-to-entropy ratio of O(1) is produced as displayed in the upper right corner of the figure (above the solid grey line). For smaller ϕ_0, the produced asymmetry will be given by

$$\frac{n_B}{s} \sim \frac{\epsilon {\phi_0}^4 {m_y}^{3/2}}{({M_X}^2 + \phi_0^2) {M_P}^{5/2} \tilde{m}} \tag{16}$$

and when $\phi_0 < \tilde{m}^{5/2} M_P^3 / m_y^{9/2}$, y's decay after ϕ's. This region is demarked by the dashed line in the Figure. Thus at $\phi = M_P$, this is the only region relevant for the dilution of the baryon asymmetry. For Planck scale baryon number violation of the type discussed below, we show the region in the $m_y - \phi_0$ plane which results in a baryon asymmetry of $10^{-10} - 10^{-11}$. Clearly a *vev* of order O(1) for ϕ results in too large an asymmetry no matter what value of m_y we choose. GUT scale baryon number violation would only compound this problem.

If we implement the constraint on m_y from their decays into LSP's [23], and consider $m_y \simeq 10^6$ GeV, and $\phi_0 = M_P$, Eq. (16) gives $n_B/s \simeq 10^{-7/2}$, and we see that the dilution is far too small. At $m_y = 10$ TeV, $n_B/s \simeq 10^{-13/2}$, the dilution is still insufficient. Even at $m_y = 100$ GeV, $n_B/s \simeq 10^{-19/2}$, the asymmetry is somewhat too large. As one can see, the primary moduli problem of entropy production is nonexistent in the AD baryogenesis scenario. Whether or not moduli can indeed be heavy enough to avoid a low reheating temperature or LSP production is another issue which has also been considered in the context of no-scale supergravity [17, 24] and recent string theory formulations [25, 26, 27].

If however R-parity is violated the constraint from an over-density in LSPs can be avoided and smaller values of m_y are allowed. For example, in orbifold compactifications of the heterotic string, the untwisted sector of the effective field theory is Heisenberg invariant at the classical level. It has recently been shown, both in a modular invariant linear multiplet formulation [28] and in a chiral multiplet formulation [29] for the dilaton supermultiplet, that dilaton stabilization with string-scale weak coupling ($g_{st} \sim 1$) can be achieved by invoking string nonperturbative effects [30]. It has further been shown [26, 27] that in these models the moduli problem of low temperature reheating can be avoided ($m_y > 10$ TeV with y a modulus t_I or the dilaton). Here we consider the quasi-realistic, modular invariant, model of [31, 26] with both gaugino and matter condensation, and moduli stabilization at a self-dual point $t_I = 1$ or $e^{i\pi/6}$ in the true vacuum. The condensation scale is

$$\Lambda_c \sim e^{-1/6b(2/g_{st}^2+\pi b_8)}, \tag{17}$$

where $3b$ is the β-function coefficient of the strong hidden sector gauge group, and $b_8 = .38 \gg b$ is the E_8 β-function coefficient. The masses of gravitino ($m_{\frac{3}{2}}$), moduli (m_t) and dilaton (m_d) in vacuum are given by

$$m_{\frac{3}{2}} = \frac{1}{4} b\Lambda_c^3 \approx \frac{b}{2b_8} m_t \approx b^2 m_d. \tag{18}$$

Thus $m_d > m_t$, and the moduli reheating problem is resolved provided $b \geq .016$–$.027$ for $g_{st} = 1$–$.5$, in which case $m_{\frac{3}{2}} \geq .74$–1.3 TeV and $\Lambda_c \geq .9 \times 10^{14}$ GeV. Scalar masses depend on their as yet unknown couplings to the Green-Schwarz term needed to restore modular invariance. If they are decoupled from this term, their masses, as specified at the condensation scale, are equal to the gravitino mass. A plausible alternative is that the untwisted sector fields couple to the GS term in the same Heisenberg invariant combination that appears in the untwisted sector Kähler potential, *i.e.* that the GS term depends only on the radii of compactification of the three tori. In this case the untwisted sector masses are $m_u(\Lambda_c) = m_t/2$. (Multi-TeV masses for some squarks and sleptons have also been proposed in other contexts [32], which are, however, subject to the stringent constraint from positivity of the scalar top mass squared [33].) If we assume that $\tilde{m} \sim m_y \sim 10$ TeV in (14), where now $\tilde{m}$ is the true vacuum mass of the particles along the relevant flat directions during inflation, one can obtain the observed baryon density of 10^{-11}–10^{-10} with values of $\phi_0 \sim 10^{-1} M_P$ if we identify M_X with the string scale $\Lambda_{st} \sim g_{st} M_P$, as would be the case in an affine level one model with no GUT below the string scale. Since now $\phi_0 < \tilde{m}^{\frac{5}{2}} M_P^3 / m^{\frac{9}{2}}$, y's decay after the ϕ's and since $m_y \lesssim 10^6$ GeV, the model falls in the shaded region of the Figure, and its viability as a mechanism for diluting baryon number requires R-parity violation. The required value of ϕ_0 can be read off the n_B/s contours in the figure at low m_y ($m_y/m_\psi \sim 10^{-8}$ in this case) after accounting for the slightly different values of M_X and $\tilde{m}$ chosen.

If R-parity is conserved, the LSP problem can be evaded only if the moduli are stabilized at or near their ground state values during inflation. In fact, the domain of attraction near $t_1 = 1$ is rather limited: $0.6 < \mathrm{Re}t_I < 1.6$, and the entropy produced by dilaton decay with an initial value in this range might be less that commonly assumed. Assuming that the entropy generated by the decay of the moduli t is negligible, the entropy generated by dilaton decay is not a problem in this model because the dilaton mass is about 10^3 TeV [26]. With untwisted sector masses ~ 10 TeV, we need $\phi_0 \leq 10^{-2} M_P$ if we include dilution from dilaton decay. Thus the amount of additional suppression needed is rather mild in these models.

One generally expects untwisted sector flat directions to be lifted by Heisenberg noninvariant terms arising from effects such as string loop threshold corrections and nonperturbative corrections, that can contain factors of the Dedekind function $\eta(t_I)$. For example, in the inflationary model of [34], some of the moduli are stabilized at $t_I = e^{i\pi/6}$; the flat directions in the corresponding untwisted sector are lifted: $m_{\phi^{AI}} \approx m_{t_I} \sim V^{\frac{1}{2}}$, and cannot contribute to the Affleck-Dine mechanism. The potential for the remaining moduli remains flat up to corrections from the condensation potential: $m_t \sim V_c^{\frac{1}{2}} \sim 100$ TeV, and the flat directions ϕ in the corresponding untwisted sector have squared masses of the same order provided $|\phi_{IA}|^2 \leq \mathrm{t_I}$. This contribution is negative if $|\phi_{IA}|^2 \leq .2\mathrm{t_I}$, and is much smaller than that induced by the loop effects ($-m_\phi^2 \sim 10^{-2}V$) in (10), and therefore cannot suppress the baryon number.

4 Regulation by Non-renormalizable Operators

Hereafter we ignore possible effects from the decay of moduli (we do however retain the dilution from inflaton decay, which is fixed and model independent provided that the inflaton decay is gravitational). We next consider the reduction in baryon asymmetry which results from the limitations on the flat direction *vev*'s which are implied by the presence in the flat direction potential of higher dimensional operators induced at the Planck scale by gravitational dynamics, or at a GUT scale (should one exist) by the presence of GUT interactions.

First we consider higher dimensional operators that are generically expected in supergravity and superstring theory; these can limit cosmological *vev*'s that are generated along flat directions of the supersymmetric standard model (including the case where the standard model is embedded in a GUT with the flat direction in question continuing to remain flat in the embedding GUT). In general, the renormalizable terms in the superpotential of the supersymmetric standard model merely represent the lowest dimensional terms in an infinite series of ascending dimension; these leading terms define the low energy Wilson effective action governing the dynamics of the standard model superfields. One expects that all terms of a given dimension that are not forbidden by gauge invariance will arise in the full effective action (global symmetries are generally violated by nonperturbative gravitational effects, and hence cannot prevent the presence of gravitationally induced terms in the effective superpotential).

On dimensional grounds one then expects that there will be a series of higher dimensional contributions to the superpotential for the flat direction field Φ, by terms of the form: $h\Phi^n/M_P^{n-3}$. These induce flat direction potential terms of the form: $h^2\phi^{(2n-2)}/M_P^{2n-6}$, where h represents a dimensionless coupling that incorporates the underlying dynamics, and we have exhibited the gravitational scale dependence explicitly. As noted above, during inflation there is a vacuum energy induced contribution to the potential of the flat directions of the form $\sim -g^2H^2\phi^2/(4\pi)^2$. Minimizing the *vev* potential, including higher dimensional corrections from a leading superpotential term of order n in Φ as above, we find that during inflation the the flat direction *vev* rolls to a value given by:

$$\phi^{(2n-4)} \sim \frac{2g^2}{(2n-2)h^2}\frac{H^2}{(4\pi)^2}M_P^{(2n-6)}. \tag{19}$$

Dropping the numerical prefactor which involves the undetermined, but generically comparable, couplings g and h, and recalling that for viable inflationary models $H/(4\pi)$ is of order $10^{-8}M_P$, we have for the flat direction *vev*'s at the end of inflation:

$$\phi^{(2n-4)} \sim (10^{-16})M_P^{(2n-4)} \tag{20}$$

In particular, for the potential arising from the leading higher dimensional gravitationally induced superpotential terms of dimension 6 ($n = 4$), we have $\phi \sim 10^{-4}M_P$.

We remark that higher dimension terms constructed from untwisted sector fields in orbifold string compactifications necessarily contain factors of the the Dedekind functions $\eta(t_I)$

that are needed to preserve modular invariance, but break Heisenberg invariance. If the superpotential contains a term $W(t) = \eta(t)^n w$, there is a corresponding contribution to the potential:

$$V(t) \ni e^K |W(t)|^2 (t+\bar{t}) \left[n^2 (t+\bar{t}) |\xi(t)|^2 - n \left(\xi(t) + \text{h.c.} \right) \right], \quad \xi = \frac{1}{\eta} \frac{\partial \eta}{\partial t}. \qquad (21)$$

Such terms have been encountered in modular invariant parameterizations of gaugino condensation [35] with $n = -6$; in that case they destabilize the potential in the direction of weak coupling (Re$s \to 0$, $K \ni -\ln(2\text{Re}s)$) because $\xi(\text{Re}t \geq 1) < 0$, and $V(t) < 0$ for $0 \leq t < 1.9$. In the present case, operators of higher dimension in untwisted fields require $n > 0$, so $V(t) \geq 0$ for Re$t \geq 1$ (it is sufficient to consider this domain because of modular invariance of the potential), and such terms are not *a priori* incompatible with a bounded potential.

As well as providing superpotential contributions to the flat direction potential, the non-perturbative gravitational effects will induce higher dimensional operators exhibiting baryon and lepton number violation, again scaled by the appropriate powers of the Planck mass. In particular, the resulting baryon and lepton number violation will, after supersymmetry breaking, give rise to quartic scalar terms in the flat direction effective potential which are now scaled by $\tilde{m}^2/M_P^2$. This results, after flat direction oscillation and decay, in a baryon asymmetry from Eq. (15) with the replacement $M_X \to M_P$

$$\frac{n_B}{s} \sim \frac{\epsilon {\phi_0}^4 {m_\psi}^{3/2}}{{M_P}^{9/2} \tilde{m}} \sim \epsilon \; 10^6 (\frac{\phi_0}{M_P})^4 \qquad (22)$$

The baryon asymmetry in this case is shown on the right hand side of the Figure ($m_y = m_\psi$) and we see that for $\epsilon \sim 1$, and ϕ_0 determined by the leading (quartic) higher dimensional gravitational contribution to the superpotential, Eq. (22) results in the observed baryon asymmetry: $n_B/s \sim 10^{-10}$. So the leading effects of quantum gravity in the minimal supersymmetric standard model act to generate baryon number violation, and potential terms limiting the flat direction *vev*'s, of precisely the magnitude needed to generate the cosmological baryon asymmetry from flat direction oscillations! This indicates how robust the mechanism is, and demonstrates that even in the absence of electroweak baryogenesis, the minimal supersymmetric standard model (including the effects of gravity) is capable of generating the BAU of the right order of magnitude. We also note that since gravity will violate $B-L$ (provided we do not gauge $B-L$; we examine GUT extensions of the gauge group below) one expects it to induce operators that violate $B-L$ (eg. through L violating operators acting on flat directions), and the resulting $B-L$ and hence B asymmetry will be immune to sphaleron erasure.

5 Regulation by GUT-scale Interactions

We now turn to a consideration of the effects on flat direction oscillations of the presence of a new large scale in addition to the Planck scale; in particular, we examine baryogenesis from flat direction oscillations in the context of Grand Unified Theories. Once one embeds the standard model in a larger gauge theory, at a high (but sub-Planckian) scale M_X, we have the possibility of both new sources of baryon number violation, and new possible contributions to the dynamics of the flat directions. In general, for all GUT structures from the minimal SU(5) gauge group on up, there will be induced baryon and lepton number violating effects in low energy processes, including baryon- and lepton-number violation inducing quartic contributions to the potential for flat direction oscillations (after supersymmetry breaking) [1].

On the other hand, a generic flat direction of the low energy supersymmetric standard model may, or may not, have its flatness lifted by GUT interactions. For those flat directions that are GUT-flat, we expect that the *vev*'s will ultimately be limited by the quantum gravity effects discussed above. If we assume that the GUT flat direction is lifted by the leading higher dimensional superpotential term (quartic) which could be induced by non-perturbative gravitational effects, then as seen above this yields a flat-direction *vev* after inflation of order $O(10^{-4})M_P$, which when substituted into the expression for the resulting GUT induced baryon asymmetry from flat direction oscillations (Eq. (15)), yields, for ϵ of order one, a baryon to entropy ratio of order 10^{-4}. If for some reason the leading gravitational superpotential operator vanished, the higher dimensional operators would stabilize the flat direction *vev* at an even larger value during inflation, resulting in an even more severe surplus of baryons.

This problematic overproduction of baryons will be regulated if the flat direction responsible for baryogenesis is flat within the supersymmetric standard model, but has its flatness lifted by GUT interactions. However, it is important to note that GUT-scale D-term interactions can never lift flat directions in the absence of supersymmetry breaking effects as long as they are F-flat. This is because of the theorem that one can always perform a complexified gauge transformation on an F-flat configuration to make it also D-flat [36]. This is counter intuitive since one may think that a flat direction can be lifted if it is no longer D-flat above the GUT-scale because of the enhanced gauge symmetry. Consider, for instance, the $L = H_u$ flat direction [4] which is D-flat under the full standard model gauge group. Since all D-flat directions in supersymmetric gauge theories are characterized by gauge-invariant polynomials [36], it is convenient to use the language of the polynomials. The flat direction under discussion corresponds to the gauge-invariant LH_u. If the gauge group is enhanced to $SO(10)$ where L belongs to a **16** and H_u to a **10**, this polynomial is no longer gauge invariant. Therefore this flat direction *stops* at the GUT-scale. However, there must be one or more Higgs fields which break the GUT gauge group to the standard model gauge group, and a certain combination of the Higgs field and the standard model flat direction remains flat. This is because one can always write a GUT-invariant polynomial using powers of the Higgs fields on top of the standard model gauge-invariant polynomial. This can be seen explicitly by studying the potential including the $U(1)_X$ D-term together with the GUT-Higgs field χ, $\bar{\chi}$ (here they are assumed to belong to $\mathbf{126}$ and $\overline{\mathbf{126}}$) with a superpotential $W = (\bar{\chi}\chi - v^2)S$:

$$V = |\bar{\chi}\chi - v^2|^2 + |S|^2(|\chi|^2 + |\bar{\chi}|^2) + \frac{g^2}{8}(|H_u|^2 - |L|^2)^2 + \frac{g'^2}{8}(2|H_u|^2 - 3|L|^2 + 10|\chi|^2 - 10|\bar{\chi}|^2)^2, \quad (23)$$

where g is the $SU(2)$ coupling and $g' = g/\sqrt{10}$ is the $U(1)_X$ coupling and S is an SO(10) singlet. $U(1)_X$ is a subgroup of SO(10) which is not contained in SU(5). One can see that along the direction $\chi = ve^{\eta}$, $\bar{\chi} = ve^{-\eta}$, $|L| = |H_u| = l$ and $l^2 = 20v^2 \sinh 2\eta$, the potential remains flat. This is because that one can construct a gauge-invariant polynomial $(\mathbf{16}\ \mathbf{10})^{10}\mathbf{126}$, and hence according to the theorem, there is a D-flat direction.

Therefore, new F-term interactions to the GUT-scale fields are necessary to lift flat directions in the standard model. For the above case, it can be the superpotential coupling **16 16 126** which makes the right-handed neutrino of the **16** heavy. The net effect then is the GUT-scale suppressed effective superpotential such as $(LH_u)(LH_u)/M_X$.

In general specific GUTs may have both GUT-flat and GUT-non-flat directions which are mutually exclusive (we describe a toy model as an example below). By mutually exclusive, we mean that the GUT-flat direction is no longer flat once the GUT-non-flat direction turns on.

As indicated above, the GUT-flat direction, with GUT scale baryon number violation leaves us with too large an asymmetry ($O(10^{-4})$. In contrast, a GUT-non-flat direction, produces a mush smaller flat direction *vev* given by eq. (20) with the replacement $M_P \to M_X$. For $M_X \sim 10^3$ and $\phi_0^4 \sim 10^{-22}$, (15) gives a baryon asymmetry of the desired magnitude. Thus, given a GUT with mutually exclusive GUT-flat and GUT-non-flat directions, unless one can a priori determine which if any is preferred, n_B/s is undetermined. In this case, a very mild form of the anthropic principle is needed: our existence indicates that a GUT-non-flat direction was chosen (perhaps randomly) over the GUT-flat one. In fact, mutual exclusivity is probably a desirable feature of the GUT with respect to the baryon asymmetry. In its absence, GUT-flat directions would remain flat despite the population of GUT-non-flat directions. In this case, once again, too large an asymmetry would be produced along the GUT-flat directions. Thus all GUT-flat directions must be mutually exclusive of some GUT-non-flat direction.

To illustrate such an example, consider a toy model with a "GUT" gauge group $U(1) \times U(1)'$, with gauge couplings $g = g'$, that is broken to $U(1)$ at the scale M_G by *vev*'s of scalar fields with only $U(1)'$ charge. We introduce the light superfields $\Phi_1, \Phi_2, \Phi_3, \Phi_4$ with $U(1) \times U(1)'$ charges $(1,1), (-1,1), (-1,-1), (2,0)$, respectively, and superpotential

$$W_L = \lambda \Phi_4 \Phi_2 \Phi_3.$$

(Anomaly cancellation in both the low energy and GUT theories is assumed to be achieved by additional $U(1)$-charged fields with no flat directions.) If $\lambda \neq 0$ there are two mutually exclusive flat directions of the low energy theory, defined by (ϕ_i is the scalar component of Φ_i) $|\phi_1| = \phi$, $|\phi_4| = 0$, and *a*) $|\phi_2| = \phi$, $|\phi_3| = 0$, or *b*) $|\phi_3| = \phi$, $|\phi_2| = 0$. If $\lambda = 0$ (*e.g.,* if there is no field Φ_4), there is a one parameter family of flat directions defined by $\phi = \phi_1 = \phi_2/\sin\alpha = \phi_3/\cos\alpha$, which connects *a*) and *b*). Above the GUT scale we introduce the massive superfields $\Pi_0, \Pi_1, \Pi_2, \Pi_3$ with $U(1) \times U(1)'$ charges $(0,0), (0,1), (0,-1), (0,-2)$, respectively, and superpotential

$$W_L = \lambda_0 \Pi_0 \left(\Pi_2 \Pi_1 - v^2 \right) + \lambda_3 \Pi_1 \Pi_3^2 + \eta \Pi_3 \Phi_2 \Phi_1.$$

(We could also add the gauge invariant term $\zeta \Pi_0 \Phi_3 \Phi_1$ but this does not affect the discussion; it just shifts the *vev* $< \pi_1 \pi_2 >$ when $\phi \neq 0$ along the flat direction *b*).) The only GUT flat direction is *b*). If $\lambda = 0$, the GUT-scale couplings can drive $\alpha \to 0$ without stabilizing ϕ.

However if $\lambda \neq 0$, and the the light fields lie along the flat direction a), ϕ can be stabilized by the GUT interactions. The effective cut-off for the low energy theory is $\Lambda = M_G$, but loops from the massive fields have the effect of shifting $\Lambda = M_G \rightarrow \Lambda = M_P$ in (10). However, integrating out these heavy fields yields higher dimension operators in the effective low energy tree potential for ϕ along a):

$$V_H = \frac{\eta^2 (\eta^2 - 2g^2) \phi^6}{2\lambda_3 v^2 + (\eta^2 - 2g^2) \phi^2} - \frac{2g^4 \phi^8}{2\lambda_0^2 v^4 - g^2 \phi^4} = \frac{\eta^2 (\eta^2 - 2g^2) \phi^6}{2\lambda_3 v^2} + O\left(\frac{\phi^2}{v^2}\right). \qquad (24)$$

Assuming $\lambda_0^2 \sim 1$, $0 < \eta^2 (\eta^2 - 2g^2) \sim 1$, $< V > \ll M_G$, we obtain a value during inflation: $\phi \sim V^{\frac{1}{4}} M_G \sim 10^{-5.5}$ for $H \sim 10^{-7}, M_G \sim 10^{-2}$.

6 Partial Sphaleron Erasure

Another possible suppression mechanism of the baryon asymmetry is to employ the smallness of the fermion masses. The baryon asymmetry is known to be wiped out if the net $B - L$ asymmetry vanishes because of the sphaleron transitions at high temperature (sphaleron erasure). However, Kuzmin, Rubakov and Shaposhnikov (KRS) [37] pointed out that this erasure can be partially circumvented if the individual $(B - 3L_i)$ asymmetries, where $i = 1, 2, 3$ refers to three generations, do not vanish even when the total asymmetry vanishes. Even though there is still a tendency that the baryon asymmetry is erased by the chemical equilibrium due to the sphaleron transitions, the finite mass of the tau lepton shifts the chemical equilibrium between B and L_3 towards the B side and leaves a finite asymmetry in the end. Their estimate is

$$B = -\frac{4}{13} \sum_i \left(L_i - \frac{1}{3}B\right) \left(1 + \frac{1}{\pi^2} \frac{m_{l_i}^2}{T^2}\right) \qquad (25)$$

where the temperature $T \sim T_C \sim 200$ GeV is when the sphaleron transition freezes out (similar to the temperature of the electroweak phase transition) and $m_\tau(T)$ is expected to be somewhat smaller than $m_\tau(0) = 1.777$ GeV. Overall, the sphaleron transition suppresses the baryon asymmetry by a factor of $\sim 10^{-6}$. This suppression factor is sufficient to keep the total baryon asymmetry at a reasonable order of magnitude in many of the cases discussed above.

Such $B-3L_i$ asymmetries between different generations can be generated by the operator

$$\tilde{m}^2 \frac{L_i^* L_j H_u^* H_u}{M_X^2}, \tag{26}$$

which violates individual $B-3L_i$ symmetries but not the total $B-L$ symmetry. Indeed, the generation of different lepton flavor asymmetries is expected within the Affleck-Dine framework. One should note however, that the KRS suppression is only viable if lepton number violating interactions for all three generations remain out-of-equilibrium. If one or more, but not all generations have lepton number violating interactions in equilibrium, then once again a large baryon asymmetry will result [38]. If all generations have interactions in equilibrium, then the asymmetry will be totally washed away by sphaleron effects.

7 Conclusions

We have considered several possibilities for controlling the potentially large baryon asymmetry produced by the Affleck-Dine mechanism in supergravity models with a Heisenberg symmetry of the Kähler potential. Moduli, which normally produce a problematic amount of entropy were shown in fact to be insufficient in diluting the baryon asymmetry in models with R-parity conservation. In the absence of R-parity, the untwisted sector of orbifold compactifications may provide a candidate for the desired dilution. Perhaps the simplest and most natural possibility for the reasonable baryon asymmetry invokes only gravitational interactions for both the baryon number violation as well as the higher dimensional operators necessary for controlling the sfermion *vevs* along flat directions. In a GUT, we showed that it is necessary to have mutually exclusive GUT-flat and GUT-non-flat directions. GUT-flat directions do not provide ample suppression of the baryon asymmetry. Finally, we noted that the mechanism suggested by KRS using lepton mass effects via sphaleron interactions could dilute an asymmetry by a factor of 10^{-6}.

Before concluding we note that string theory predicts a relation between the values of the gauge coupling g at unification and the scale of unification, $\Lambda_G \sim gM_P$, that is not satisfied in fits of the data to the MSSM. It has recently been realized that in the strongly coupled limit of heterotic superstring compactifications ("M-theory"), string unification may in fact be effectively 5-dimensional [39, 40], with the extra dimension "turning on" at $O(10^{14})$ GeV for the gravitational modes of the theory. In this scenario the two E_8's of the heterotic string live on two ten-dimensional surfaces that are separated by this fifth (or eleventh) dimension r. This changes the dimensional analysis that relates the Plank scale to the string tension, and modifies the above relation in such a way that consistency with MSSM RGE's can be achieved for the choice $r \sim 10^4$. One might expect that $r^{-1} \sim 10^{-4}$ would represent an upper limit for values of *vev*'s that may be reliably considered in the low energy effective theory. However, just as the MSSM gauge coupling RGE's are left intact, because the observed sector is contained in a (weakly coupled) E_8 that does not see the extra dimension, the AD flat directions that are charged under the MSSM gauge group will also be unaffected. Unless effective interactions at the scale of the fifth dimension can be generated along the flat directions in such a way so as to preserve the running of the gauge couplings, these effects will not be capable of reducing the baryon asymmetry.

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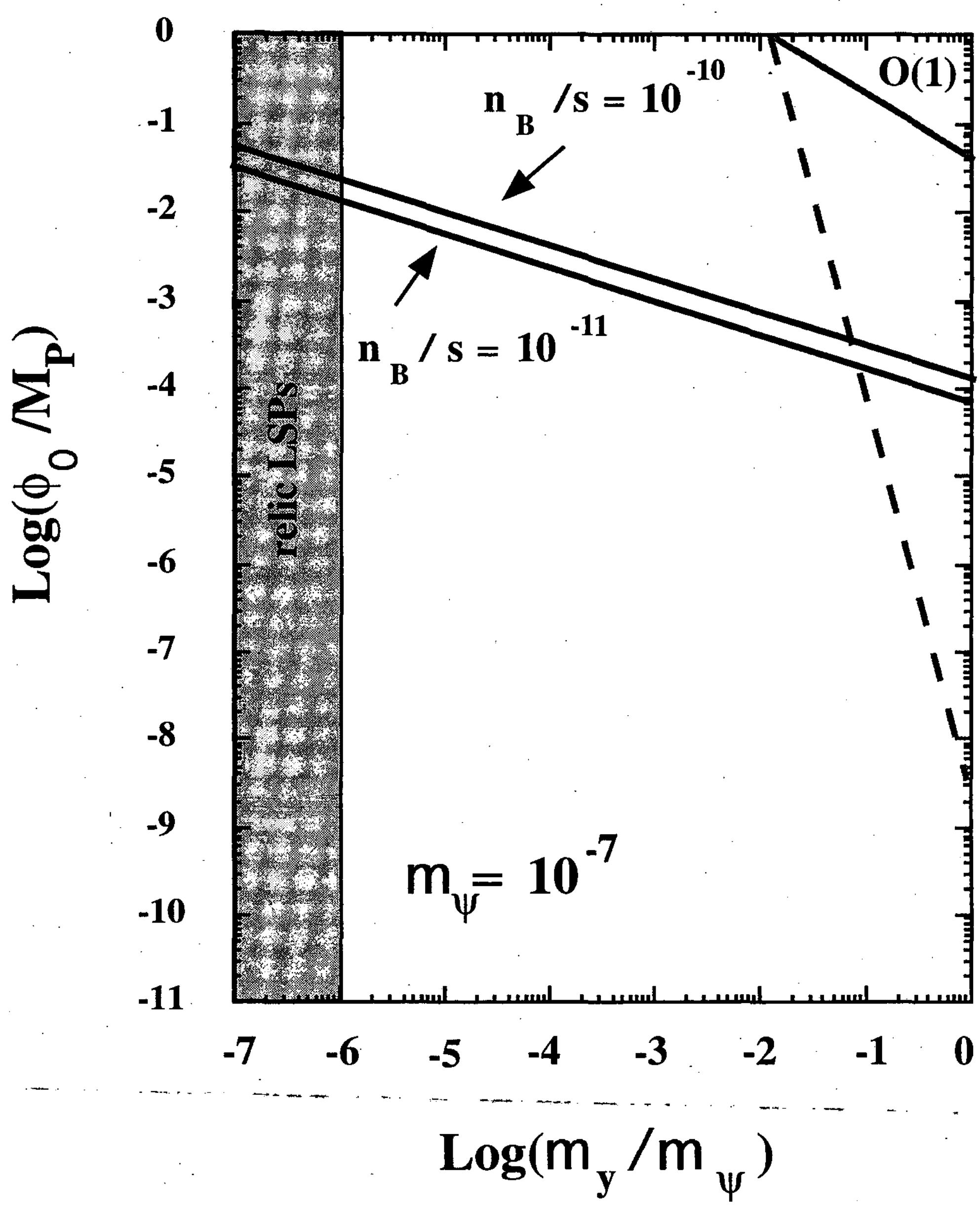
Figure Captions

Figure 1: Features of the baryon asymmetry in the $\phi_0 - m_y$ plane. In the upper right hand corner of the plot (above the light grey line), the baryon asymmetry is $n_B/s \sim O(1)$ and is due to the late decays of sfermion oscillations which have come to dominate the energy density of the Universe. Between the grey solid and dashed lines, the sfermions decay earlier, while the Universe is dominated by the radiation products of either y or ψ decay (depending on whether m_y is less than m_ψ or not. Below the dashed line the sfermion flat directions decay before the y's or ψ's. Below the grey line, the baryon asymmetry is given by eq. 16 unless $m_y > m_\psi$ in which case it is given by eq. 15. Also shown on the figure (in black) are lines of constant $n_B/s = 10^{-10}$ and 10^{-11}. $m_\psi = 10^{-7}$, with the Hubble parameter during inflation, $H_I \sim m_\psi$, as well as $\tilde{m} = 10^{-16}$ were chosen. In this case, Planck scale baryon number violation ($M_X \sim 1$) was assumed. Finally, the shaded region to the left is excluded if R-parity is conserved.



O(1)
n B /s = 10 -10
n B / s = 10 -11
relic LSPs
m ψ = 10 -7
Log(φ 0 /M P)
Log(m y /m ψ)
0
-1
-2
-3
-4
-5
-6
-7
-8
-9
-10
-11
-7
-6
-5
-4
-3
-2
-1
0

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